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NONLINEAR THEORY OF THE COLLISIONAL RAYLEIGH-TAYLOR INSTABILITY--ETC(U)

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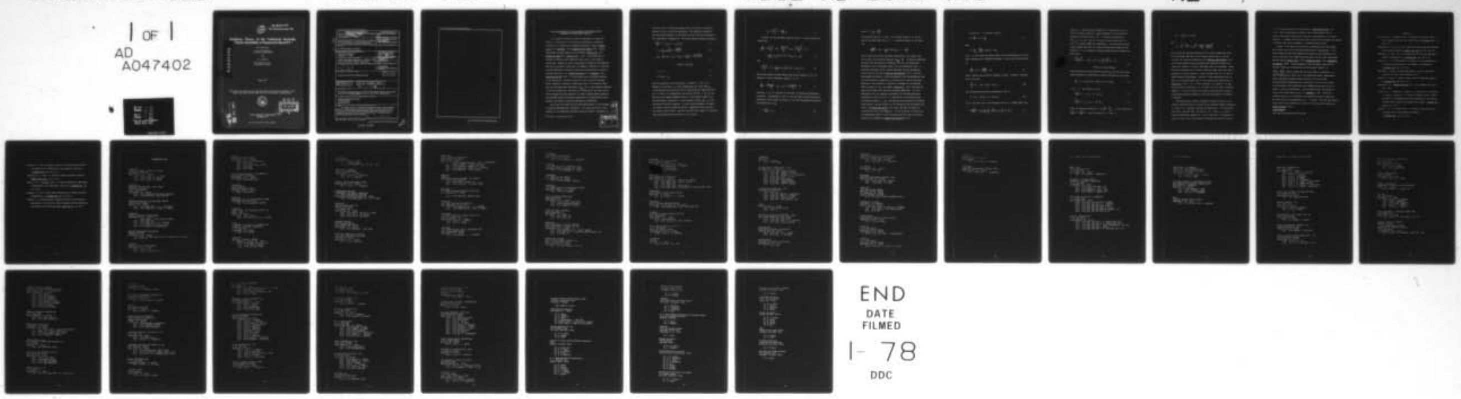

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Nonlinear Theory of the Collisional Rayleigh-Taylor Instability in Equatorial Spread F

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The nonlinear behavior of the collisional Rayleigh-Taylor instability is studied in equatorial Spread F by including a dominant two-dimensional nonlinearity. It is found that on account of this nonlinearity, the instability saturates by generating damped higher spatial harmonics. This saturated power spectrum for the density fluctuations goes like k^{-2}. A comparison between experimental observations and theory is presented.

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NONLINEAR THEORY OF THE COLLISIONAL RAYLEIGH-TAYLOR INSTABILITY IN EQUATORIAL SPREAD F

It is generally believed now that the phenomenon of equatorial Spread F is initiated by a collisional Rayleigh-Taylor instability occuring in the evening time bottomside equatorial F region [Balsley et al. 1972; Haerendel, 1974; Scannapieco and Ossakow 1976]. Recent experimental evidence supports this postulate [Kelley et al. 1976; Woodman and LaHoz 1976; McClure et al., 1977]. This instability is believed to generate large transverse scale sizes, on the order of several kms. Whereas at shorter scalelengths, a variety of other physical effects are believed to take place, for example, transition to collisionless Rayleigh-Taylor excitation, drift wave excitation due to linear or nonlinear processes, etc. [Hudson and Kennel, 1975; Haerendel, 1974; Chaturvedi and Kaw, 1976]. In the present paper, we confine ourselves to the question of nonlinear saturation of the collisional Rayleigh-Taylor instability occuring at large scale sizes. In so doing, we find that the two-dimensional nonlinearity $(\underline{v}_{\underline{E}\times\underline{B}} \cdot \nabla\tilde{n})$ in the continuity equation is the most dominant one, and leads to a saturation of the linear instability through nonlinear excitation of linearly damped spatial harmonics. The saturated density fluctuations amplitudes are presented and are discussed in the light of available experimental observations.

We now wish to present the two-dimensional analysis of the nonlinear collisional Rayleigh-Taylor instability. For simplicity we use a slab geometry for the nighttime bottomside equatorial Spread F ionosphere, in

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which the z-axis is along the magnetic field, the x-axis is vertically upwards, and the y-axis points eastwards. The background ionospheric electron density gradient is thus along the x-axis, and the acceleration due to earth's gravity is antiparallel to it. The two-fluid equations describing the system are

$$\frac{\partial n_\alpha}{\partial t} + \nabla \cdot (n_\alpha \underline{v}_\alpha) = - \nu_R (n_\alpha - n_{\alpha o}) \tag{1}$$

$$0 = -T_e \nabla n_e - e n_e \left(-\nabla \phi + \frac{\underline{v}_e \times \underline{B}_o}{c} \right) \tag{2}$$

$$m_i n_i \left(\frac{\partial}{\partial t} + \underline{v}_i \cdot \nabla \right) \underline{v}_i = -T_i \nabla n_i + e n_i \left(-\nabla \phi + \frac{\underline{v}_i \times \underline{B}_o}{c} \right) + m_i n_i \underline{g} - m_i n_i \nu_{in} \underline{v}_i \tag{3}$$

$$\nabla \cdot \underline{J} = 0 \tag{4}$$

$$\underline{J} = ne(\underline{v}_i - \underline{v}_e) \tag{5}$$

where the subscript α denotes species e (electron), i (ion); n is density; $\underline{v}$ is velocity; ν_R is the recombination rate (note that $n_{\alpha o}$ is the equilibrium density such that in equilibrium the right hand side of (1) is zero); T is temperature; the electric field $\underline{E} = - \nabla \phi$; m is the mass; $\underline{B}_o$ is the ambient magnetic field, c is the speed of light; g is the acceleration of gravity; e is the electronic charge; and ν_{in} is the ion-neutral collision frequency. In the following treatment, we ignore ion inertia (valid for large scale sizes), temperature effects and any zero-order electric fields. Then the ions drift with respect to electrons, with an equilibrium velocity given by $(\Omega_i \equiv eB_o/m_i c)$

$$v_{\underline{i}o} = \frac{\underline{g} \times \hat{z}}{\Omega_i} + \frac{\nu_{in}}{\Omega_i^2}\underline{g} \tag{6}$$

Further, the ion continuity equation and Eq. (4) can be written as, respectively

$$\frac{\partial \tilde{n}}{\partial t} + \frac{\underline{g} \times \hat{z}}{\Omega_i} \cdot \nabla\tilde{n} - \frac{c\nabla\tilde{\phi} \times \hat{z}}{B_o} \cdot \nabla n_o - \frac{c\nabla\tilde{\phi} \times \hat{z}}{B_o} \cdot \nabla\tilde{n}$$

$$- \frac{c}{B_o\Omega_i}\left[n\nu_{in}\nabla^2\tilde{\phi} + \nabla\tilde{\phi} \cdot \nabla(\tilde{n}\nu_{in})\right] = -\nu_R\tilde{n} \tag{7}$$

and

$$\frac{\underline{g} \times \hat{z}}{\Omega_i} \cdot \nabla\tilde{n} - \frac{c}{B_o\Omega_i}\left[n\nu_{in}\nabla^2\tilde{\phi} + \nabla\tilde{\phi} \cdot \nabla(\tilde{n}\nu_{in})\right] = 0 \tag{8}$$

where tilde denotes the perturbation and we have assumed $\tilde{n}_e \simeq \tilde{n}_i \simeq \tilde{n}$. Equation (7) can be rewritten, using Eq. (8), as

$$\frac{\partial \tilde{n}}{\partial t} - \frac{c\nabla\tilde{\phi} \times \hat{z}}{B_o} \cdot \nabla n_o + \nu_R\tilde{n} = \frac{c\nabla\tilde{\phi} \times \hat{z}}{B_o} \cdot \nabla\tilde{n} \tag{9}$$

Equations (8) and (9) describe the nonlinear collisional Rayleigh-Taylor instability. Linearizing (8) and (9), the set readily yields an absolute instability with a growth rate given by (for waves propagating horizontally and perpendicular to $\underline{B}_o$)

$$\gamma = \frac{g}{L\nu_{in}} - \nu_R \tag{10}$$

where $L^{-1} = \frac{1}{n_o} \frac{dn_o}{dx}$.

The nonlinear terms in (8) and (9) are the last terms in (8) and the term on the right hand side of (9). A comparison between the two shows that

$$\frac{c\nabla\tilde{\phi}\times\hat{z}}{B_o} \cdot \nabla\tilde{n} : \frac{c}{B_o\Omega_i} \nabla\tilde{\phi}\cdot\nabla(\tilde{n}\nu_{in}) \approx 1 : \frac{\nu_{in}}{\Omega_i}$$

Thus we find that the most important nonlinearity in the present case is the one in the continuity equation, $[\tilde{\underline{v}}_{\underline{E\times B}} \cdot \nabla\tilde{n}]$. It should be mentioned here that this nonlinearity is important only for a two-dimensional perturbation and is vanishingly small for a one-dimensional case. An elaborate discussion of the relative importance of this nonlinearity can be found in the work of Rognlein and Weinstock [1974]. They have investigated the effects of this nonlinearity in the two-dimensional case for the cross-field instability in the equatorial E region electrojet. Physically, the cross-field instability is analogous to the Rayleigh-Taylor instability (see, for example, Sudan et al., 1973) and hence the similarity between their nonlinear saturation mechanisms. In a one-dimensional case, the important nonlinearity for the Rayleigh-Taylor instability is the one in Eq. (8) (note that in a collisionless case, the Pedersen effects, $\sim \nu_{in}/\Omega_i$, are replaced by polarization ion inertial drift effects), and the evolution of the instability can be markedly different (see Chaturvedi and Kaw, 1975, where they have discussed the collisionless case in one-dimension). In the following, we discuss the two-dimensional effects for the collisional case only and our analysis parallels the treatment of Rognlein and Weinstock [1974].

We treat Eq. (8) linearly, and write

$$\tilde{\psi} \equiv \frac{e\tilde{\phi}}{T} = -i\beta \frac{\tilde{n}}{n_o} \tag{11}$$

where

$$\beta = \frac{g}{\nu_{in}} \frac{k_y \Omega_i}{k^2 c_s^2} \quad \text{and} \quad c_s^2 = T/m_i$$

In Eq. (9), the second term gives growth of the perturbations, while the third represents the recombination damping. We can rewrite this equation as

$$\frac{\partial \tilde{n}}{\partial t} = \gamma \tilde{n} + \frac{c\nabla\tilde{\phi} \times \hat{z}}{B_o} \cdot \nabla \tilde{n} \tag{12}$$

where γ denotes the net growth (damping) of modes. Consider a perturbation of the form

$$\frac{\tilde{n}_1}{n_o} = A_{1,1} \sin (k_y y - \omega t) \cos k_x x \tag{13a}$$

and the associated potential perturbation, from (11)

$$\tilde{\psi}_1 = \beta A_{1,1} \cos(k_y y - \omega t) \cos k_x x \tag{13b}$$

Use of (13a) and (13b) in the nonlinear term of (12) readily shows that

$$\frac{c\nabla\tilde{\phi} \times \hat{z}}{B_o} \cdot \nabla \tilde{n} = n_o \left(\frac{g}{2\nu_{in}} \frac{k_y^2}{k^2} k_x A_{1,1}^2 \right) \sin 2k_x x \tag{14}$$

Equation (14) shows that the interaction of a mode given by (13) on itself nonlinearly generates a spatial harmonic in the x-direction through the nonlinear term in (12). We denote this harmonic by $(\tilde{n}/n_o)_{2,0} = A_{2,0} \sin\ 2k_x x$, and it can be verified easily from (12) that it is linearly damped by recombination. The interaction of this linearly damped mode with the linearly growing mode through the nonlinearity in (12), introduces a nonlinear damping for the linearly unstable mode,

$$\frac{c\nabla\tilde{\phi}_{1,1} \times \hat{z}}{B_o} \cdot \nabla\tilde{n}_{2,0} = n_o\left(-2\, k_x \frac{g}{\nu_{in}} \frac{k_y^2}{k^2} A_{1,1}\, A_{2,0}\right) \cdot \sin(k_y y - \omega t) \cos k_x x \cos 2k_x x \tag{15}$$

These interactions can be written down explicitly in the form of mode coupled equations by substituting a general perturbation of the form

$$\frac{\tilde{n}}{n_o} = A_{1,1} \sin(k_y y - \omega t) \cos k_x x + A_{2,0} \sin 2k_x x \tag{16}$$

in Eq. (12). Upon doing so, we get

$$\frac{\partial A_{1,1}}{\partial t} = \gamma_{1,1} A_{1,1} - 2\alpha A_{1,1} A_{2,0} \tag{17}$$

$$\frac{\partial A_{2,0}}{\partial t} = \gamma_{2,0} A_{2,0} + \frac{\alpha}{2} A_{1,1}^2 \tag{18}$$

where the coupling coefficient $\alpha = k_x \frac{k_y^2}{k^2} \frac{g}{\nu_{in}}$. In the steady-state, $\frac{\partial A_{1,0}}{\partial t} = \frac{\partial A_{2,0}}{\partial t} = 0$, and we obtain from (17) and (18):

$$A_{2,0} = \frac{\gamma_{1,1}}{2\alpha} \approx 1/(2k_x L) \tag{19}$$

and

$$A_{1,1} = \left(- \frac{2\gamma_{2,0}}{\alpha} A_{2,0}\right)^{\frac{1}{2}} \approx \left(\frac{k^2}{k_y^2} \frac{\nu_R \nu_{in}}{gLk_x^2}\right)^{\frac{1}{2}} \tag{20}$$

We note that the saturated amplitude of the linearly damped mode turns out to be larger than that of the linearly growing mode, as was the case for the cross-field instability too [Rognlein and Weinstock, 1974]. In the above analysis we have considered the interaction between two modes only; however, in reality, many modes should be considered. Hence, the actual saturation of the linear instability might take place after a generation of several harmonics. Thus, estimates given by (19) and (20) might represent overestimates. However, a crude estimate may still be made of the saturated amplitudes for the coherent wave observations of Dyson et al. [1974]. At ~ 425 kms. altitude, if we assume L ≃ 30 kms, then the density fluctuation is ≈ 0.5% which is on the order of observed magnitude. An initial problem with several modes requires solving the coupled mode equations numerically, and such investigation is currently in progress.

The above analysis clearly represents a coherent evolution of the primary instability which, on account of the nonlinear harmonic generation, is likely to be a steepened structure observationally. Such structures would have a spectral power law going as $\sim k^{-2}$ (see Eq. 19), which is what the observations suggest too. On the other hand, a turbulent generation of drift waves by the primary irregularities would also lead to

a k^{-2} law obeying spectrum of irregularities [Chaturvedi and Kaw, 1975; 1976]. Thus an experimental distinction can be made between the two kinds of conditions by doing phase correlation measurements on the irregularities. A coherent evolution would show a phase-correlated spectrum; whereas, a turbulent case would correspond to a randomly-phased spectra.

Finally, a few words on the subject of irregularities on the topside and a lower limit on the scale sizes from physical considerations. It is widely agreed now that the irregularities growing on the bottomside, convect up (and down) due to $\underline{E} \times \underline{B}$ drifts associated with the local polarization electric fields [Kelley et al., 1976; Woodman and LaHoz, 1976; Scannapieco and Ossakow, 1976]. A crude estimate of these irregular convective velocities (which can be associated with the Spread F bubbles) is $\sim\left(\frac{g}{\nu_{in}}\frac{\delta n}{n_o}\right)$, where δn represents the local density fluctuation. As these irregularities convect up to the topside and create irregularities there, they generate smaller scale sizes due to the well known two-step process, through the drift instability. Now it is well known that the kinetic drift instability has a maximum growth rate at $k_\perp \rho_i \geqslant 1$ (where ρ_i is the ion gyroradius), and nonlinearly also, the spectra of kinetic drift waves peaks at $k_\perp \rho_i \geqslant 1$ [Kadomtsev, 1965]. Hence, the minimum scale size, to be observed in the equatorial Spread F region may be of the order of ion gyroradius, i.e., ~ 3 meters. This may qualitatively explain the cut-off observed at 3 meters in the irregularity spectrum by Woodman and Basu [1977].

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